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ELECTRON VELOCITY SHEAR INSTABILITY IN THE AURORAL IONOSPHERE.(U)
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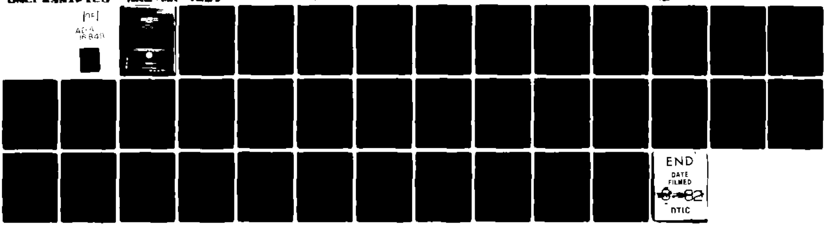

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A linear electrostatic kinetic theory of velocity sheared electron streams flowing parallel to the geomagnetic field is presented. Our development, which includes the effects of density gradients, background ions and electron collisions, is valid for arbitrary T_e/T_i, where T_e and T_i are the electron and ion temperatures, respectively. For applications to the auroral ionosphere, we find that the interface between downward flowing hot electrons and upward flowing return current cold electrons, commonly seen near auroral arcs, is unstable to a high frequency electron velocity shear driven instability. The mode is characterized by $\gamma \lesssim 10^{-2}\,\omega_{pe}$ and $r_{Li} > k^{-1} > r_{Le}$, where γ is the growth rate, ω_{pe} is the electron

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20. ABSTRACT (Continued)

plasma frequency, k is the wavenumber, and $r_{L\alpha}$ is the mean Larmor radius of species α. In the non-linear regime, it is anticipated that this instability can evolve into a vortex configuration and as a result, can act as a block near the edges of the electron flow.

CONTENTS

ELECTRON VELOCITY SHEAR INSTABILITY IN THE AURORAL IONOSPHERE

I. Introduction

Electron streams flowing in magnetic fields are a common occurrence in terrestrial and astrophysical plasmas, e.g., the earth's ionosphere and magnetosphere [Arnoldy, 1974· Anderson and Vondrak, 1975], and solar flares [Sturrock, 1974]. These streams have been observed flowing parallel or perpendicular to the magnetic field with the flow often spatially varying perpendicular to the magnetic field. It is well know that this flow configuration is unstable [Chandrasekhar, 1957] and subject to a Kelvin-Helmholtz-like velocity shear instability. For the ionosphere-magnetospheric plasma, Webster [1957] suggested the possible connection between small scale structures in auroral arcs and instabilities arising from electron $\underline{E} \times \underline{B}$ flows perpendicular to the magnetic field. However, there is considerable evidence [Arnoldy 1981, and references therein] that strong inhomogeneous electron flows parallel to the geomagnetic field also exist near auroral arcs. To our knowledge, the instabilities of transverse sheared electron streams flowing parallel to the magnetic field has not been fully discussed in a space plasma physics context.

In the following we present a linear kinetic theory of transverse velocity sheared electron flows parallel to the magnetic field (Harrison, 1963; Harrison and Stringer, 1963; Mikhailovskii and Rukhadze, 1966 Rome and Briggs, 1972). Our treatment, which includes density gradients, arbitrary T_e/T_i, background ions and electron collisions, is applicable to the low β (particle pressure/magnetic pressure) weakly ionized auroral ionospheric plasma. We show that the interface between downward flowing hot electrons and upward flowing return current cold electrons is unstable to a high frequency electron velocity shear driven instability. In Section II we derive a general

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kinetic theory and present our principal results. In Section III we summarize our findings and discuss the nonlinear implications of this instability.

II. Theory

A. Physical Mechanism of the Instability

We first present a simple physical picture of the instability driven by a transversely sheared electron flow along a magnetic field (Rome and Briggs, 1972). The magnetic field $\underset{\sim}{B} = B\,\hat{e}_z$ is taken to be constant in space and time. The ions are assumed to form an immobile, homogeneous background, while the electrons are assumed to have an inhomogeneous flow velocity $\underset{\sim}{V}_{oe} = V_{oe}(x)\hat{e}_z$ such that $V_{oe}(x) = (x-x_o)\,\partial V_{oe}/\partial x$ with $\partial V_{oe}/\partial x < 0$. We impose upon this configuration a two-dimensional, perturbed electric field $\delta E = \delta E_y\hat{e}_y + \delta E_z\hat{e}_z$ as shown in Fig. (1). The perturbation force in the z-direction acting on an electron fluid element at $x = x_o$, can be written

$$\delta F_{ez} = \delta F^a_{ez} + \delta F^d_{ez}$$

where

$$\delta F^a_{ez} = -\,e\,\delta E_z$$

$$\delta F^d_{ez} = -\,m_e\,\delta V_{ex}(\partial V_{oe}/\partial x)$$

and $\delta V_{ex} = -\,c\delta E_y/B$ is the perturbed $\delta\underset{\sim}{E} \times \underset{\sim}{B}$ drift due to δE_y. Here, δF^a_{ez} is the "acceleration" force due to δE_z which acts to neutralize the charge imbalance, while δF^d_{ez} is the "deceleration" force due to the convection term $(\delta\underset{\sim}{V}\cdot\nabla)\underset{\sim}{V}$ and opposes δF^a_{ez} when $\partial V_{oe}/\partial x < 0$. Thus, when $\delta F^d_{ez} > \delta F^a_{ez}$, the force imbalance on an electron fluid element leads to an enhancement of charge separation, and hence, instability. For electrostatic fluctuations,

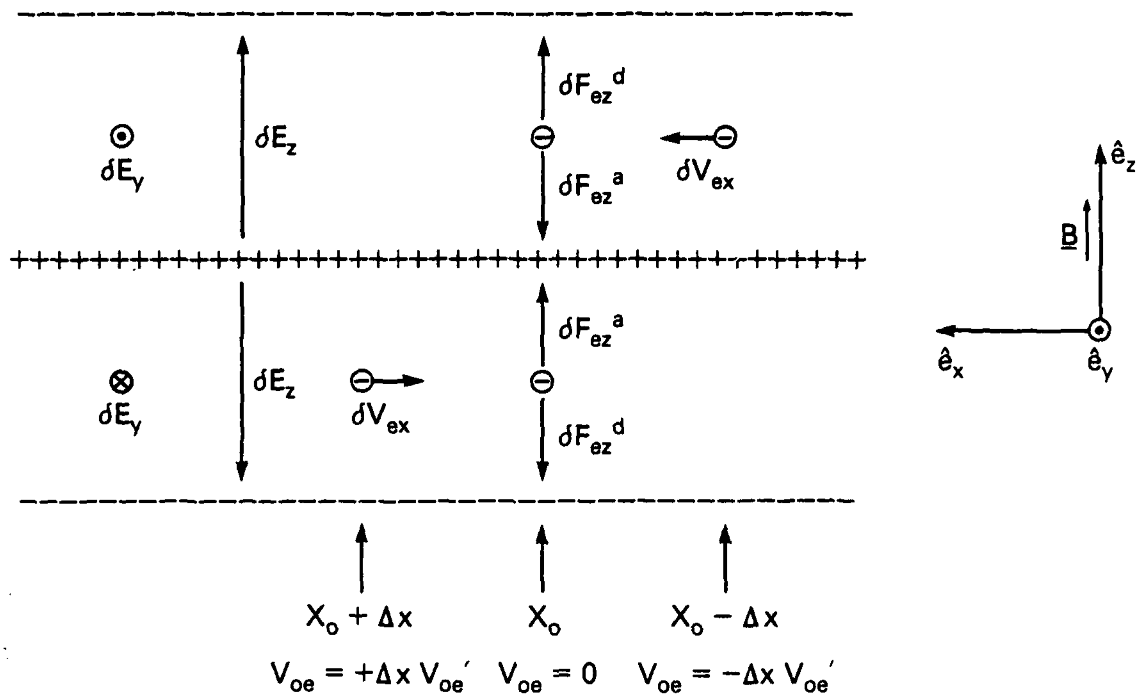

Fig. 1 — Schematic drawing illustrating basic physical mechanism of electron velocity shear instability. The xz-plane lies in the plane of the paper with the y-axis directed out of the paper. The positive (+) signs represent excess positive charge while the negative (–) signs indicate negative charge.

$\delta E_z = - i\, k_z \delta\phi$ and $\delta E_y = - i\, k_y\, \delta\phi$, where $\delta\phi$ is the electrostatic potential; this gives the following instability criterion $|(1/\Omega_e)\partial V_{eo}/\partial x| > |k_z/k_y|$. We prove this criterion rigorously in the next sections.

B. Dispersion Equation

We consider a plasma configuration as shown in Figure 2. Note that the ions are assumed to have no flow in the z-direction and that the temperature is taken to be constant. We consider only electrostatic oscillations and assume perturbed quantities vary as $\exp[i\,(k_y y+k_z z-\omega t)]$ with $k_z^2/k_y^2 << 1$. We assume a weakly inhomogeneous plasma ($r_{Li}^2 << L^2$ where r_{Li} is the mean ion Larmor radius and L is the scale length of the density and velocity gradients) and make use of the local approximation ($k_y^2 >> \partial^2/\partial x^2 >> 1/L^2$). We assume that $\Omega_i^2 << \omega^2 << \Omega_e^2$ so that the ions can be considered unmagnetized and the electrons magnetized. Finally, we include electron collisions in the analysis, i.e., electron-ion collisions.

The unperturbed electron distribution function is given by

$$F_{oe}(v_\perp^2,\, v_\parallel,\, X) = n_{oe}(X)(\pi v_e^2)^{-3/2} \exp[-(v_\perp^2+(v_\parallel-V_{oe}(X))^2)/v_e^2] \qquad (1)$$

where $v_\perp^2 = v_x^2 + v_y^2$, $v_\parallel = v_z$ and $X = x - v_y/\Omega_e$ are constants of motion. Here, $v_e = (2T_e/m_e)^{1/2}$ and $\Omega_e = eB_o/m_e c$ are the electron thermal velocity and cyclotron frequency, respectively. We expand (1) about x = 0 and obtain

$$F_{oe}(v_\perp^2,\, v_\parallel,\, X) \simeq F_{oe}^m\,(v_\perp^2,\, v_\parallel) - \frac{v_y}{\Omega_e}\, \frac{\partial F_{oe}}{\partial x}\bigg|_{x=0} \qquad (2)$$

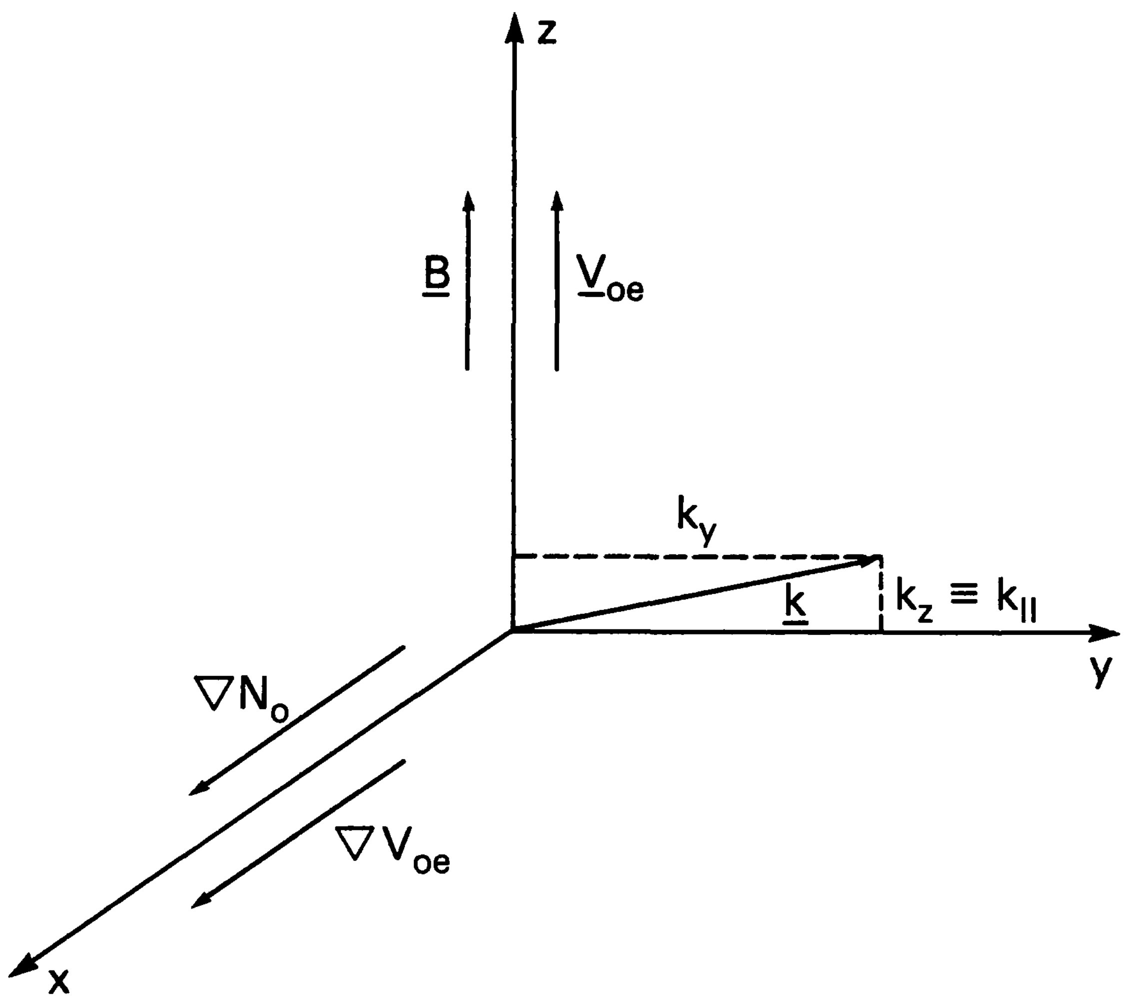

Fig. 2 — Basic cartesian geometry used to discuss electron velocity shear instability. The k-vector shown lies in the yz-plane. The electron velocity, magnetic field, and z-axis are collinear.

where

$$F_{oe}^{m}\ (v_\perp^2,\ v_\parallel) = n_{oe}(\pi v_e^2)^{-3/2}\ \exp[-(v_\perp^2+(v_\parallel - V_{oe})^2)/v_e^2]\ \Big|_{\ x\ =\ o} \tag{3}$$

and

$$\frac{\partial\ F_{oe}}{\partial x}\Bigg|_{\ x=0} = F_{oe}^{m}\ \left[\frac{\partial\ \ell n\ n_{oe}}{\partial x} + \frac{2(v_\parallel - V_{oe})V_{oe}}{v_e^2}\ \frac{\partial\ \ell n\ V_{oe}}{\partial x}\right]_{x=0} \tag{4}$$

The equilibrium electron drifts associated with (2) are $V_{ox}^{e} = 0$, $V_{oy}^{e} = -\ (v_e^2/\partial\Omega_e)\partial\ \ell n\ n_{oe}/\partial x$ and $V_{oz}^{e} = V_{oe}$ where $\underset{\sim}{V}_o^e = \int d^3v\ \underset{\sim}{v}\ F_{oe}$. The electron equilibrium pressure tensor is defined as $\underset{\approx}{P}_{oe} = \int d^3v\ (\underset{\sim}{v} - \underset{\sim}{V}_o{}^e)(\underset{\sim}{v} - \underset{\sim}{V}_o{}^e)\ F_{oe}$ and is given by

$$\underset{\approx}{P}_{oe} = n_{oe}T_e \begin{pmatrix} 1 & 0 & 0 \\ 0 & 1 & A_e \\ 0 & A_e & 1 \end{pmatrix} \tag{5}$$

where $A_e = -\ (1/\Omega_e)\partial V_{oe}/\partial x$. It is the anisotropy in the pressure tensor that provides the free energy to drive the instability.

The unperturbed ion distribution function is assumed to be

$$F_{oi} = n_{oi}(\pi v_i^2)^{-3/2}\ \exp\ [-(v_x^2 + (v_y - V_{di})^2 + v_z^2)/v_i^2] \tag{6}$$

where $v_i = (2T_i/m_i)^{1/2}$ is the ion thermal velocity, $V_{di} = (v_i^2/2\Omega_i)\ \partial\ \ell n\ n_{oi}/\partial x$ is the ion diamagnetic drift velocity *and* $\Omega_i = eB_o/m_i c$ *is the ion clclotron* frequency.

The electrostatic dispersion equation is obtained from Poisson's equation

$$\nabla^2\delta\phi = 4\pi(\delta n_e - \delta n_i) \tag{7}$$

where $\delta\phi$ is the perturbed potential, δn_σ is the perturbed density of species σ, i.e, $\delta n_\sigma = \int d^3v\, \delta f_\sigma$, and δf_σ is the perturbed distribution function of species σ. We define $\chi_\sigma \delta\phi = -(4\pi e_\sigma/k^2)\delta n_\sigma$ so that the dispersion equation is

$$D(\omega,k) = 1 + \chi_i + \chi_e = 0 \tag{8}$$

We obtain χ_e by making use of the BGK or Krook model to describe the effects of electron collisions (Kadomtsev, 1965). That is, we consider

$$\left(\frac{\partial}{\partial t} + \underset{\sim}{v} \cdot \frac{\partial}{\partial \underset{\sim}{x}} - \frac{e}{m_e}\frac{1}{c}\underset{\sim}{v} \times \underset{\sim}{B} \cdot \frac{\partial}{\partial \underset{\sim}{v}}\right)\delta f_e = -\frac{e}{m_e}\, i\underset{\sim}{k}\cdot \frac{\partial F_{oe}}{\partial \underset{\sim}{v}}\,\delta\phi - \nu_e\left(\delta f_e - \frac{\delta n_e}{n_o} F_{oe}\right). \tag{9}$$

where ν_e represents the electron collision frequency. Following standard techniques, i.e., the method of characteristics, we determine δf_e from Eqs. (2) and (9); then integrating over velocity space we finally obtain

$$\chi_e = \frac{2\omega_{pe}^2}{k^2 v_e^2}\left[1+\left[\left(\zeta_e - \frac{k_y}{k_z}\frac{V_{de}}{v_e}\right)Z(\zeta_e) - \frac{1}{2}\frac{k_y}{k_z}\frac{1}{\Omega_e}\frac{dV_{oe}}{dx} Z'(\zeta_e)\right]\Gamma_o(b_e)\right] \times$$
$$\left[1+\frac{i\nu_e}{k_z v_e} Z(\zeta_e)\, \Gamma_o(b_e)\right]^{-1} \tag{10}$$

where $\zeta_e = (\omega + i\nu_e - k_z V_{oe})/k_z v_e$, $\omega_{pe}^2 = 4\pi n e^2/m_e$, $\Gamma_o(b_e) = \exp(-b_e)\, I_o(b_e)$, $b_e = k_y^2 r_{Le}^2/2$, $r_{Le} = v_e/\Omega_e$, I_o is the modified Bessel function of order 0, $V_{de} = -(v_e^2/2\Omega_e)\, \partial \ln n/\partial x$, Z is the plasma dispersion function and $Z'(\zeta) = dZ/d\zeta$. The ion response χ_i is simply given by

$$\chi_i = \frac{2\omega_{pi}^2}{k^2 v_i^2}\left[1 + \zeta_i\, Z(\zeta_i)\right] \tag{11}$$

where $\omega_{pi}^2 = 4\pi ne^2/m_i$, $\zeta_i = (\omega - k_y V_{di})/kv_i$ and $V_{di} = (v_i^2/2\Omega_i)\partial \ln n/\partial x$. Thus, Eqs. (8), (10) and (11) describe electrostatic waves in the presence of a sheared parallel electron velocity for the frequency regime $\Omega_i^2 \ll \omega^2 \ll \Omega_e^2$.

C. Analytical Analysis

In order to gain insight into the nature of the electron velocity shear instability, we analytically solve Eq. (8) in two limiting cases: "cold" electrons ($T_e \ll T_i$) and "hot" electron ($T_e \gg T_i$). For simplicity, we expand about $x = x_o$ and an consider electron velocity profile such that $V_{oe} = (x-x_o)\,(\partial V_{oe}/\partial x)_{x=x_o}$ so that $V_{oe}(x=x_o) = 0$.

1. Cold electrons ($T_e \ll T_i$)

In this limit, the contribution from the ions can be ignored since $\zeta_i \ll 1$ and the dispersion equation can be written

$$D(\omega,k) = 1 + \chi_e = 0 \tag{12}$$

In order to simplify Eq. (10) we make the following assumptions: $(\omega + i\,\nu_e)/k_z v_e \ll 1$ and $k_y^2 r_{Le}^2 \ll 1$. For these conditions we note that $Z(\zeta_e) \simeq -1/\zeta_e - 1/2\zeta_e^2$, $Z'(\zeta_e) = 1/\zeta_e^2$ and $\Gamma_o(b_e) \simeq 1 - k_y^2 r_{Le}^2/2$. We find that

$$\chi_e = \left[\frac{\omega_{pe}^2}{\Omega_e^2} - \frac{\omega_{pe}^2}{(\omega+i\nu_e)^2}\frac{k_z}{k_y}\left(\frac{k_z}{k_y} + \frac{1}{\Omega_e}\frac{dV_{oe}}{\partial x}\right) - \frac{\omega_{pe}}{\omega+i\nu_e}\frac{\omega_{pe}}{\Omega_e}\frac{1}{k}\frac{\partial \ln n}{\partial x}\right]\left(1 - \frac{i\,\nu_e}{\omega+i\nu_e}\right)^{-1} \tag{13}$$

Making use of Eqs. (12) and (13) we arrive at the following equation

$$D(\omega,k) = (1+\sigma^2)\tilde{\omega}^2 - (i\tilde{\nu}_e - \sigma/k_y L_n)\tilde{\omega} - \Lambda = 0 \tag{14}$$

where $\tilde{\omega} = (\omega+i\nu_e)/\omega_{pe}$, $\tilde{\nu}_e = \nu_e/\omega_{pe}$, $\sigma = \omega_{pe}/\Omega_e$, $L_n = (\partial\ \ell n\ n/\partial x)^{-1}$, $\Lambda = (k_z/k_y)(k_z/k_y + V'_{oe}/\Omega_e)$, and $V'_{oe} = \partial V_{oe}/\partial x$. The solution of Eq. (14) is

$$\tilde{\omega} = \frac{1}{2}\ [(i\nu_e - \sigma/k_y L_n) \pm \{(i\nu_e - \sigma/k_y L_n)^2 + 4\Lambda(1+\sigma^2)\}^{1/2}](1+\sigma^2)^{-1} \tag{15}$$

We further simplify Eq. (15) by examining the following limits.

a. Collisionless plasma ($\nu_e = 0$)

In the collisionless limit ($\nu_e = 0$), Eq. (15) reduces to

$$\omega = \frac{1}{2}\ [-\frac{\sigma}{k_y L_n} \pm \{\frac{\sigma^2}{k_y^2 L_n^2} + 4\Lambda(1+\sigma^2)\}^{1/2}]\ (1 + \sigma^2)^{-1} \tag{16}$$

Instability can occur when

$$\Lambda = \frac{k_z}{k_y}\ (\frac{k_z}{k_y} + \frac{1}{\Omega_e}\ \frac{\partial V_{oe}}{\partial x}) < -\frac{1}{4}\ \frac{\sigma^2}{1+\sigma^2}\ \frac{1}{k_y^2 L_n^2} \tag{17}$$

In the limit $L_n \to \infty$, this instability criterion is identical with that derived from the simple physical picture presented in Sec. IIA. The important aspect of Eq. (17) is that the density gradient is a stabilizing influence (Mikhailovskii and Rukhadze, 1966). Moreover, the density gradient is most effective in stabilizing the instability when $\sigma^2 \gg 1$ (i.e., $\omega_{pe}^2 \gg \Omega_e^2$) and $k_y L_n \ll 1$. The latter condition indicates that the density gradient acts to stabilize long wavelength modes before short wavelength modes, so that a "long wavelength cut off" should exist in inhomogeneous plasmas. We also note that the instability attains maximum growth for

$$\frac{k_z}{k_y} = -2\frac{1}{\Omega_e}\frac{\partial V_{oe}}{\partial x} \tag{18}$$

b. Collisional plasma ($\nu_e \neq 0$)

We consider a collisional plasma ($\nu_e \neq 0$), but neglect the density gradient for simplicity ($L_n \to 0$). In this limit, Eq. (15) becomes

$$\tilde{\omega} = i\frac{\tilde{\nu}_e}{2}(1+\alpha^2)^{-1}\left[1 \pm \left(1 - \frac{4\Lambda(1+\alpha^2)}{\tilde{\nu}_e^2}\right)^{1/2}\right] \tag{19}$$

A necessary condition for instability is

$$\Lambda = \frac{k_z}{k_y}\left(\frac{k_z}{k_y} + \frac{1}{\Omega_e}\frac{\partial V_{oe}}{\partial x}\right) < 0 \tag{20}$$

which is the same condition as Eq (17) for $L_n \to \infty$. The growth rate is given by

$$\tilde{\gamma} = \left(\frac{|\Lambda|}{1+\alpha^2}\right)^{1/2} - \frac{\tilde{\nu}_e}{2}\frac{1+2\alpha^2}{1+\alpha^2} \quad \text{for } \tilde{\nu}_e^2 \ll 4\Lambda(1+\alpha^2) \tag{21}$$

and

$$\tilde{\gamma} = \frac{|\Lambda|}{\tilde{\nu}_e} - \frac{\alpha^2}{1+\alpha^2}\tilde{\nu}_e \quad \text{for } \tilde{\nu}_e^2 \gg 4\Lambda(1+\alpha^2) \tag{22}$$

Equation (21) is the weakly collisional limit and Eq. (22) is the strongly collisional limit. A key parameter in determining the strength of the instability in a collisional plasma is $\alpha = \omega_{pe}/\Omega_e$. In fact, for $\alpha^2 \gg 1$ it is found that collisions prevent the mode from becoming unstable. However, for $\alpha^2 \to 0$, instability presists, even for a large collision frequency.

2. Hot electrons ($T_e >> T_i$)

In this limit the contribution from the ions cannot be ignored and the dispersion equation is

$$D(\omega,k) = 1 + \chi_i + \chi_e \tag{23}$$

Again, to simplify Eq. (23) and to focus on the role of the ions, we make the following simplifying assumptions: $\nu_e = 0$, $L_n \to \infty$, $\omega >> k_z v_e$, $k_y^2 r_{Le}^2 << 1$ and $\omega >> k_z v_i$. The susceptibilities χ_σ are now given by

$$\chi_e = \frac{\omega_{pe}^2}{\Omega_e^2} - \frac{\omega_{pe}^2}{\omega^2} \frac{k_z}{k_y} \left(\frac{k_z}{k_y} + \frac{1}{\Omega_e} \frac{\partial V_{oe}}{\partial x}\right) \tag{24}$$

and

$$\chi_i = - \frac{\omega_{pi}^2}{\omega^2} \tag{25}$$

so that the dispersion equation is

$$1 + \frac{\omega_{pe}^2}{\Omega_e^2} - \frac{\omega_{pe}^2}{\omega^2} \left(\frac{k_z^2}{k_y^2} + \frac{k_z}{k_y} \frac{1}{\Omega_e} \frac{\partial V_{oe}}{\partial x} + \frac{m_e}{m_i}\right) = 0 \tag{26}$$

or

$$(1 + \alpha^2)\, \tilde{\omega}^2 - (\Lambda + \frac{m_e}{m_i}) = 0 \tag{27}$$

Clearly, instability can occur when

$$\Lambda = \frac{k_z}{k_y} \left(\frac{k_z}{k_y} + \frac{1}{\Omega_e} \frac{\partial V_{oe}}{\partial x}\right) < - \frac{m_e}{m_i} \tag{28}$$

Increased stability is found when finite electron temperature is taken into account [Harrison, 1963]. The influence of the ions is to make the instability criterion more stringent than the cold electron case, i.e., $\Lambda < 0$. Physically, this is due to the fact that the ions are no longer immobile and are able to respond to the perturbation field δE_z (see Fig. 1) Also, Eq. (28) indicates that the instability is easier to excite in "heavy ion" plasmas.

However, we can show that an instability can be excited even when Eq. (28) is not satisfied, i.e , $\Lambda > - m_e/m_i$. In this case it is crucial that electron resonance terms be retained in Eq. (24), which have been neglected. We expand the electron Z function as $Z(\zeta_e) = - 1/\zeta_e - 1/\zeta_e^2 + i\pi^{1/2} \exp(-\zeta_e^2)$ and $Z'(\zeta_e) = 1/\zeta_e^2 - i\pi^{1/2}\, 2\zeta_e \exp(-\zeta_e^2)$ so that dispersion equation becomes

$$D(\omega,k) = 1 + \alpha^2 - \frac{1}{\tilde{\omega}^2}\left(\Lambda + \frac{m_e}{m_i}\right) + i\pi^{1/2} \frac{2\omega_{pe}^2}{k_z^2 v_e^2} \zeta_e \exp(-\zeta_e^2)\, \Lambda = 0 \tag{30}$$

where $\zeta_e = \omega/k_z v_e$. We assume $\Lambda + m_e/m_i > 0$ so that

$$\omega_r = \omega_{\ell h}\left[1 + \frac{m_i}{m_e}\Lambda\right]^{1/2} \tag{31}$$

where $\omega_{\ell h} = \omega_{pi}/(1 + \alpha^2)^{1/2}$. Assuming $\gamma << \omega_r$, it is easily shown that

$$\gamma = - \pi^{1/2} \frac{\omega_{pe}^2}{k_z^2 v_e^2} \zeta_e \exp(-\zeta_e^2)\, \Lambda\, (1+\alpha^2)^{-1}\, \omega_r \tag{32}$$

so that instablity can occur when $\Lambda < 0$; the same instability criterion as in the cold electron case. An important aspect of this instability is that it has a real frequency associated with it (Eq. (31)), while the other instabilities discussed are purely growing (i.e., $\omega_r = 0$) in the absence of a density gradient. We add that, as in the cold electron limit, both a density gradient and electron collisions have a stabilizing influence on this resonant electron velocity shear instability.

D. Numerical Results

To better illustrate the various limits presented in the previous subsection, we present a set of curves which solve Eq. (8) exactly for a broad range of parameters. The important parameters that are varied are T_e/T_i, $k_y r_{Le}$, r_{Le}/L_n, $V_o/\Omega_e L$, $(\partial V_{oe}/\partial x = V_o \ /L)$, ν_e/Ω_e and ω_{pe}/Ω_e. We also consider an O^+ plasma for these cases, i.e., $m_i = 16\ m_p$. For applications to the auroral ionosphere we use parameters typical of regions in and near discrete auroral arcs. The transverse velocity shear for electron flow parallel to the geomagnetic field can be approximated by $|\partial V_{oe}/\partial x| \approx |V_c/L|$ where L is a typical latitudinal arc width and $V_o \approx v_h - v_c$ is the difference in parallel streaming velocities between the hot downward flowing current carrying electrons (v_h) and the upward flowing cold return current electrons (v_c). Typical values for the latitudinal width of discrete arcs are in the range $L \approx 100$ m – 10 km [Davis, 1978]. For v_c we take $j_{\parallel,c} \approx en_c v_c \approx 10^{-6}$ A/m^2 [Anderson and Vondrak, 1975 and references therein] with $n_c \approx 10^4 - 10^5$ cm^{-3} giving $v_c \approx 10^2 - 10^3$ cm/sec. For v_h we consider the field-aligned downward flowing electrons with energies $E \approx 1 - 10$ keV [Arnoldy, 1981] giving $v_h \approx 10^9$ cm/sec. As a result, an estimate for the transverse shear is $|\partial V_{oe}/\partial x| \approx |(v_h - v_c)/L| \approx 10^4$ sec^{-1} with L = 1 km giving for the parameter $|V_o/\Omega_e L| \approx 10^{-2}$ where $\Omega_e \approx 10^6$ sec^{-1}.

In Fig. 3 we plot ω/ω_{pe} vs. T_e/T_i. The solid lines show the growth rate γ/ω_{pe} while the dashed lines show the real frequency ω_r/ω_{pe}. The parameters used are $k_y r_{Le} = 0.1$, $\omega_{pe}/\Omega_e = 0.5$, $\nu_e/\omega_{pe} = 0$,

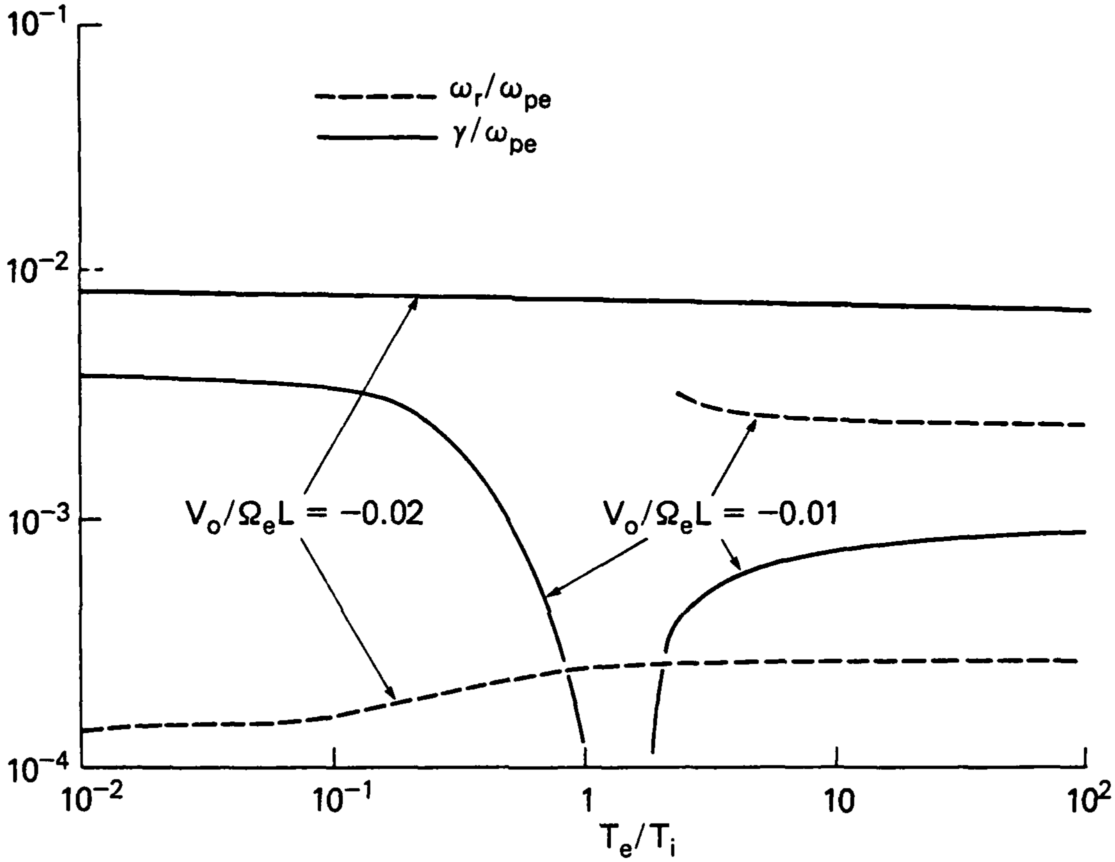

Fig. 3 — Plot of ω/ω_{pe} vs. T_e/T_i for strong shear ($V_o/\Omega_e L = -0.02$) and weaker shear ($V_o/\Omega_e L = -0.01$). Note stabilization for $T_e/T_i \simeq 1$ for weak shear case.

$r_{Le}/L_n = 1.0 \times 10^{-4}$ for $V_o/\Omega_e L = -0.01$ we choose $k_z/k_y = 0.005$, and for $V_o/\Omega_e L = -0.02$ we choose $k_z/k_y = 0.010$.

We first discuss the strong shear case, $V_o/\Omega_e L = -0.02$. We note that the instability criteria $\Lambda = (k_z/k_y)(k_z/k_y + V_o/\Omega_e L) < 0$ and $\Lambda = (k_z/k_y)(k_z/k_y + V_o/\Omega_e L) < -m_e/m_i$ are satisfied for the parameters chosen. As expected, a nonresonant instability occurs for the entire range of T_e/T_i considered. Note that the growth rate is slightly smaller for $T_e/T_i = 10^2$ than for $T_e/T_i = 10^{-2}$. This is due to the stabilizing influence of the nonresonant ion contribution (Eq. (25)). Also, there is a small real frequency associated with this mode, $\omega_r/\omega_{pe} \lesssim 2 \times 10^{-4}$, which is due to the density gradient.

The weak shear case, $V_o/\Omega_e L = -0.01$, is chosen such that $\Lambda = (k_z/k_y)(k_z/k_y + V_o/\Omega_e L) < 0$ but $\Lambda = (k_z/k_y)(k_z/k_y + V_o/\Omega_e L) > -m_e/m_i$. We anticipate a nonresonant instability in the "cold" electron regime ($T_e \ll T_i$) and a resonant instability in the "hot" electron regime ($T_e \gg T_i$). This is clearly illlustrated in Fig. 3. In the limit $T_e \ll T_i$, a strong nonresonant instability occurs which asymptotes to a growth rate of $\gamma/\omega_{pe} \simeq 4.6 \times 10^{-3}$ in the limit $T_e \to 0$. Also, there is a small real frequency associated with this mode due to the density gradient which, for the most part, has $\omega_r/\omega_{pe} < 1.0 \times 10^{-4}$ for $T_e < T_i$. As $T_e/T_i \to 1$, the ions become important and stabilize the mode because of ion Landau damping, i.e., $\omega \sim kv_i$. On the other hand, for $T_e \gg T_i$ the resonant instability described in Sec. II. B. 2 is shown. Associated with this mode is a large real frequency (Eq. (31)). The growth rate is comparable to, but smaller, than the nonresonant "cold" electron mode. Again, as $T_e/T_i \to 1$, the assumptions used to arrive at Eq. (32) break down and the mode is stabilized by ion Landau damping.

In Fig. 4 we plot γ_m/ω_{pe} vs. $k_y r_{Le}$ for $V_o/\Omega_e L = -0.01$, $\omega_{pe}/\Omega_e = 0.5$, $T_e/T_i = 10^{-3}$ and 10^3, and $r_{Le}/L_n = 0$ and 1.0×10^{-4}. Here, γ_m denotes the growth rate maximized with respect to k_z/k_y; typically, $k_z/k_y \lesssim 0.005$ for the curves shown. We first discuss the "cold" electron case ($T_e/T_i = 10^{-3}$). Instability extends over a very broad range in k_y space from $k_y r_{Le} \approx 10^{-3}$ to $k_y r_{Le} \approx 1$. The growth rate maximizes for $k_y r_{Le} \approx 0.1$. As $k_y r_{Le} \to 0$, the instability is stabilized because of ion Landau damping. The additional damping due to a density gradient is evident by contrasting the $r_{Le}/L_n = 0$ and $r_{Le}/L_n = 1.0 \times 10^{-4}$ curves in the long wavelength limit ($k_y r_{Le} \ll 1$). The mode is stabilized in the short wavelength limit ($k_y r_{Le} > 1$) because of finite electron Larmor radius effects.

The situation for the resonant instability in the "hot" electron regime ($T_e/T_i = 10^3$) is somewhat different In this case the waves are much more localized in k_y space, with growth mostly occurring in the narrow region $k_y r_{Le} \approx 0.1 - 1.0$. The growth rate is sharply reduced as $k_y r_{Le}$ becomes small (i.e., $k_y r_{Le} < 0.1$) because the resonant electron contribution, proportional to $\exp(-\omega_r^2/k_z^2 v_e^2)$, becomes negligible in this limit. Again, for $k_y r_{Le} > 1$, the mode is stabilized because of finite electron Larmor radius effects.

In Fig. 5 we plot γ/ω_{pe} vs. ν_e/ω_{pe} for $k_y r_{Le} = 0.1$, $k_z/k_y = 0.005$, $r_{Le}/L_n = 1.0 \times 10^{-4}$, $V_o/\Omega_e L = 0.01$, $T_e/T_i = 10^{-3}$ and 10^3, and $\omega_{pe}/\Omega_e = 0.5$ and 2.0. The important features of this figure are the following. First, electron collisions eventually stabilize the instability regardless of T_e/T_i. Second, electron collisions are much more effective at stabilizing the mode when ω_{pe}/Ω_e is large, which is consistent with the analytical analysis presented in Sec. II. B. 1. Finally, there appears to be a weak collisional instability for the "hot" electron case ($T_e/T_i = 10^3$) when $\omega_{pe}/\Omega_e > 1$.

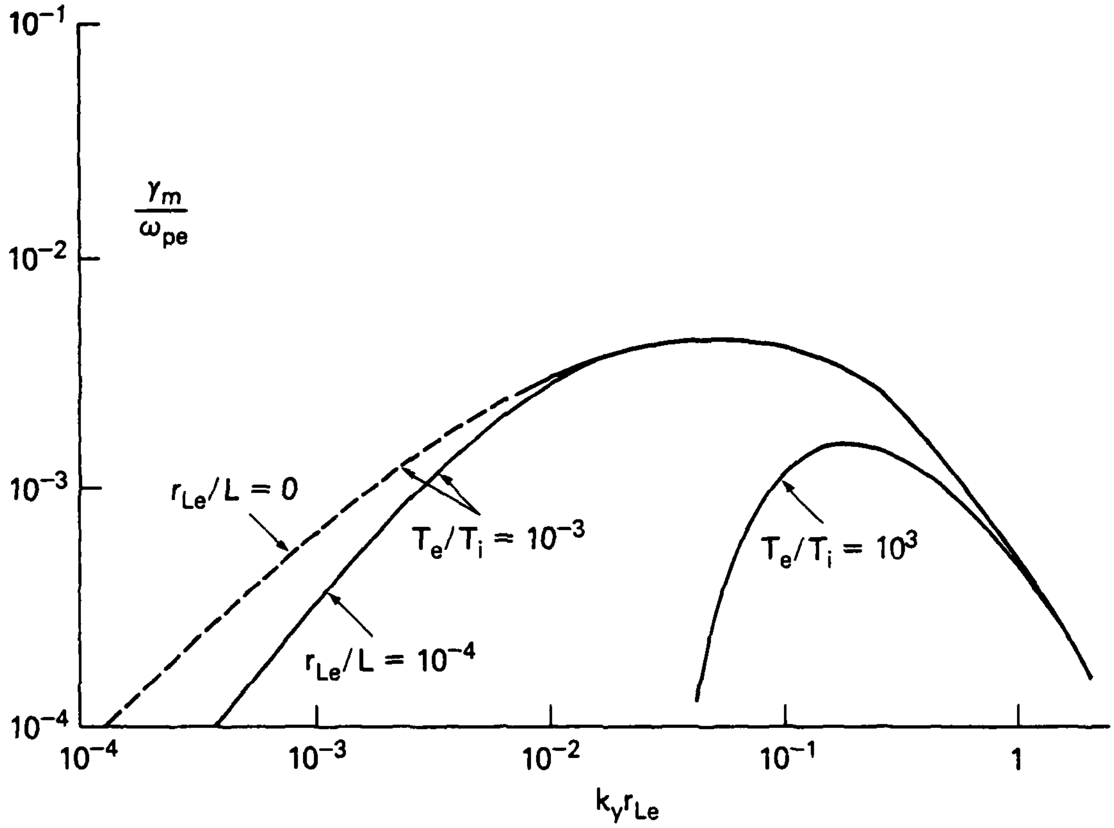

Fig. 4 — Plot of γ_m/ω_{pe} vs, $k_y r_{Le}$ for $V_o/\Omega_e L = -0.01$, $\omega_{pe}/\Omega_e = 0.5$, $T_e/T_i = 10^{-3}$, 10^3, and $r_{Le}/L_n = 0$, 1×10^{-4}. Note narrowing of region of unstable waves for $T_e/T_i = 10^3$.

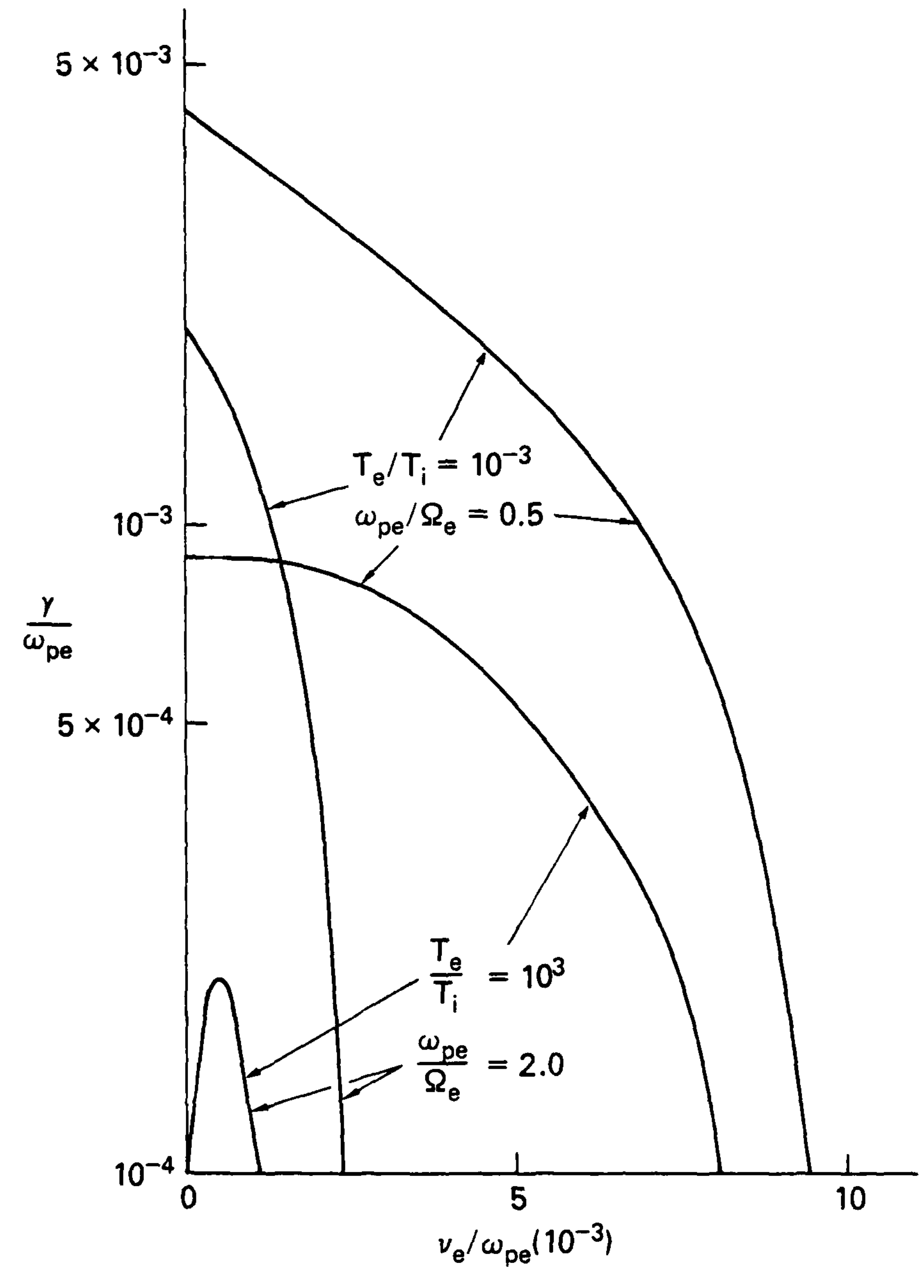

Fig. 5 — Graph of γ/ω_{pe} vs. ν_e/ω_{pe} for $k_y r_{Le} = 0.1$, $k_z/k_y = 0.005$, $r_{Le}/L_n = 10^{-4}$, $V_o/\Omega_e L = -0.01$, $T_e/T_i = 10^{-3}$, 10^3 $\omega_{pe}/\Omega_e = 0.5$, 2. Note reduction in growth rate for increasing collision frequency.

III. Summary and Discussion

We have presented a linear, electrostatic, kinetic theory of velocity sheared electron streams flowing parallel to a magnetic field. For simplicity we have treated the case where the electron velocity varies transverse to the direction of the flow. Since a Vlasov analysis has been used, finite Larmor radius effects and wave-particle resonances have been properly treated. These are important since (1) we find appreciable wave growth for $kr_{Le} \lesssim 1$, (2) electron-wave resonances can be a destabilizing influence, and (3) ion-wave resonances are a stabilizing influence. In addition, we have included an electron density gradient transverse to the flow direction, background ion response, and electron collisional effects. Moreover, through numerical solution of the fundamental dispersion relation, our theory is valid for arbitrary values of T_e/T_i. Our principal results can be summarized as follows:

1. For "cold" ($T_e << T_i$) velocity sheared electron streams, we find a nonresonant instability when

$$\frac{k_z}{k_y}\left(\frac{k_z}{k_y} + \frac{1}{\Omega_e}\frac{\partial V_{oe}}{\partial x}\right) < 0 \qquad (33)$$

   We note that a density gradient is stabilizing and provides a long wavelength cutoff such that the unstable waves are preferentially excited locally ($k^2L^2 > 1$).

2. For "hot" ($T_e >> T_i$) velocity sheared electron streams, our results again indicate a nonresonant instability when

$$\frac{k_z}{k_y}\left(\frac{k_z}{k_y} + \frac{1}{\Omega_e}\frac{\partial V_{oe}}{\partial x}\right) < -\frac{m_e}{m_i} \qquad (34)$$

As in the cold electron case, the density gradient exerts a stabilizing influence. However, we have also found a resonant (electron-wave) instability that has the same instability condition as the cold electron case (Eq. (33)).

3. For $T_e \approx T_i$, except for strong shears, (Eq. (34)), the modes are stabilized due to ion Landau damping. Ion Landau damping also stabilizes the long wavelength modes (i.e., $k_y r_{Le} \to 0$) for $T_e/T_i \neq 0$.

4. Electron collisions exert a stabilizing influence independent of T_e/T_i. Collisions are more effective for stabilization when $\omega_{pe}/\Omega_e > 1$.

In the preceeding theoretical development we have assumed that the unstable modes have high frequencies $\gamma > \Omega_i$ such that the background ions cannot execute a gyro oscillation on the instability time scale. Thus, the ions are assumed unmagnetized. For $\gamma \lesssim \Omega_i$, the ions must be considered magnetized and the previous analysis is invalid. In this case, other instabilities, e.g., cyclotron instabilities [Kindel and Kennel, 1972], are possible. We have shown a posteriori that the condition $\gamma > \Omega_i$ has been met for the parameter regimes studied.

Nonlinear numerical studies of velocity sheared ions flowing parallel to a magnetic field have recently been performed [Tajima and Leboeuf, 1980]. These simulations indicate classical vortex formation and the production of anomalous viscosity resulting in a quasilinear flattening of the initially sheared velocity profile. We anticipate similar behavior for sheared electron streams and defer a detailed study of the nonlinear evolution of electron velocity shear instabilities in space plasmas to a later report.

For application of this theory to the auroral ionosphere, we consider the interface, usually seen near discrete arcs, between downward flowing high velocity hot electrons and upward streaming low velocity cold return current electrons. Many experimental studies of auroral arc structure and dynamics have been made, particularly rocket and satellite measurements. Kelley and Carlson [1977] have detected near the edge of an F-region auroral arc, electrostatic waves with spatial scale sizes less than the measured ambient velocity shear scale size. These results are not inconsistent with our findings of a local ($k^2 L^2 >> 1$) electrostatic instability due to transverse sheared electron streams flowing parallel to the geomagnetic field. In addition, several other investigators [Whalen and McDiarmid, 1972; Bryant et al., 1973; Arnoldy et al., 1974] have found that high energy field-aligned electrons are generally localized near the edges of auroral arcs in regions separating different plasmas. These electrons have also been found to occur in bursts. Moreover, electron energy spectra near the edges of auroral arcs [Bryant et al., 1973 Bryant, 1981 Carlson and Kelley, 1977] often show structure, i.e., cold low energy electron are intermixed with hot high energy electrons and vice versa. These features, observed near the edges of auroral arcs, may be explained by the linear and nonlinear evolution of the parallel electron velocity shear instability. This can be seen by the following scenario. The parallel electron velocity shear instability can evolve, in the nonlinear regime, into vortices which can be described as a turbulent boundary layer near the edges of the electron flow. The vortices by their very nature, will tend to mix and transport hot low density electrons with cold high density thermal electrons. These vortical structures will convert kinetic streaming energy into vortical rotational energy and thereby act as a block near the edges of the electron flow. This blockage can be described as a "resistivity" and lead to the formation of localized electric fields and subsequent acceleration. As the vortices grow and expand, the initial velocity profile will flatten in a quasilinear fashion due to anomalous viscosity effects [Tajima and Leboeuf, 1980; Miura and Sato, 1978]. As a results this "resistivity" will be sporadic and bursty in character as the velocity sheared profile alternately steepens and flattens. These hypotheses will be tested in detail in future work.

Since, for conditions typical of discrete F-region auroral arcs, i.e. $L = 1\text{-}10$ km, $k_z/k_y = 10^{-2}$, $V_{oe}/\Omega_e L = 10^{-2}$, we find unstable wavelengths ranging ranging from $\lambda = 1 - 100$ m, these irregularities could be observed using radar backscatter methods [Hanuise et al., 1981]. We note that since collisions are stabilizing, this electron velocity shear instability will be preferentially excited at high altitudes. In addition, throughout our theoretical treatment of this instability, we have ignored beam instabilities [Dungey and Strangeway 1976; Papadopoulous et al., 1974] and their interactions with velocity gradient driven modes. Finally, these results may be applicable to the more diffuse interface between region 1 and 2 current systems in the auroral ionosphere [Iijima and Potemra, 1976] and to electron current return current regions in solar flares [Knight and Sturrock, 1977].

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